

UT-16-24  
IPMU-16-0082  
TU-1024  
June, 2016

# Upper Bound on the Gluino Mass in Supersymmetric Models with Extra Matters

Takeo Moroi<sup>(a,b)</sup>, Tsutomu T. Yanagida<sup>(b)</sup> and Norimi Yokozaki<sup>(c)</sup>

<sup>(a)</sup>*Department of Physics, University of Tokyo, Tokyo 113-0033, Japan*

<sup>(b)</sup>*Kavli Institute for the Physics and Mathematics of the Universe (Kavli IPMU),  
University of Tokyo, Kashiwa 277-8583, Japan*

<sup>(c)</sup>*Department of Physics, Tohoku University, Sendai 980-8578, Japan*

## Abstract

We discuss the upper bound on the gluino mass in supersymmetric models with vector-like extra matters. In order to realize the observed Higgs mass of 125 GeV, the gluino mass is bounded from above in supersymmetric models. With the existence of the vector-like extra matters at around TeV, we show that such an upper bound on the gluino mass is significantly reduced compared to the case of minimal supersymmetric standard model. This is due to the fact that radiatively generated stop masses as well the stop trilinear coupling are enhanced in the presence of the vector-like multiplets. In a wide range of parameter space of the model with extra matters, particularly with sizable  $\tan\beta$  (which is the ratio of the vacuum expectation values of the two Higgs bosons), the gluino is required to be lighter than  $\sim 3$  TeV, which is likely to be within the reach of forthcoming LHC experiment.

# 1 Introduction

Although the low-energy supersymmetry (SUSY) is attractive from the points of view of, for example, naturalness, gauge coupling unification, dark matter, and so on, to which the standard model (SM) has no clue, no signal of the SUSY particles has been observed yet. Thus, one of the important questions in the study of models with low-energy SUSY is the scale of SUSY particles.

It is well-known that the observed Higgs mass of  $\sim 125$  GeV [1] gives information about the mass scale of SUSY particles (in particular, stops). The Higgs mass is enhanced by radiative corrections when the stop masses are much larger than the electroweak scale [2–6]. Thus, the stop masses are bounded from above in order not to push up the Higgs mass too much; the stop masses are required to be smaller than  $10^4 - 10^5$  GeV as far as  $\tan\beta$ , which is the ratio of the vacuum expectation value of the up-type Higgs to that of the down-type Higgs, is larger than a few. (For the recent study of such an upper bound, see, for example, [7].) Then, too large gluino masses are also disfavored because, via renormalization group (RG) effects, it results in stops which are too heavy to make the Higgs mass consistent with the observed value. Such an upper bound on the gluino mass is important for the future collider experiments, in particular, for the LHC Run-2, in order to discover and to study models with low energy SUSY. The purpose of this letter is to investigate how such an upper bound on the gluino mass depends on the particle content of the model.

We pay particular attention to SUSY models with extra vector-like chiral multiplets which have SM gauge quantum numbers. In these days, such extra vector-like matters are particularly motivated from the excess of the diphoton events observed by the LHC [8–11]. The most popular idea to explain the diphoton excess is to introduce a scalar boson  $\Phi$  with which the LHC diphoton excess can be due to the process  $gg \rightarrow \Phi \rightarrow \gamma\gamma$ . In such a class of scenarios, vector-like particles which interact with  $\Phi$  are necessary to make  $\Phi$  being coupled to the SM gauge bosons. Indeed, it has been shown that the LHC diphoton excess are well explained in SUSY models with vector-like chiral multiplets [12–30]. Assuming a perturbative gauge coupling unification at the GUT scale of  $\sim 10^{16}$  GeV,<sup>#1</sup> three or four copies of the vector-like multiplets, which transform  $\mathbf{5}$  and  $\bar{\mathbf{5}}$  in  $SU(5)$  gauge group, are suggested, and their masses need to be around or less than 1 TeV. In addition, the vector-like chiral multiplets are also motivated in models with non-anomalous discrete  $R$ -symmetry [32, 33].

With extra vector-like chiral multiplets, the RG evolutions of the coupling constants and mass parameters of the SUSY models drastically change compared to those in the minimal SUSY SM (MSSM). Consequently, as we will see, the upper bound on the gluino mass becomes significantly reduced if there exist extra vector-like chiral multiplets. Such an effect has been discussed in gaugino mediation model [34] and in the light of recent diphoton excess at the LHC [26].

In this letter, we study the upper bound on the gluino mass in SUSY models with extra vector-like matters, assuming more general framework of SUSY breaking. We extend the

---

<sup>#1</sup>For the perturbativity bounds on models with extra matters, see [31].

previous analysis and derive the upper bound on the gluino mass. We will show that the bound on the gluino mass is generically reduced with the addition of extra matters. The upper bound becomes lower as the number of extra matters increases, and the bound can be as low as a few TeV which is within the reach of the LHC Run-2 experiment.

## 2 Enhanced Higgs boson mass and gluino mass

We first explain how the upper bound on the gluino mass is reduced in models with extra vector-like multiplets. To make our discussion concrete, we consider models with extra chiral multiplets which can be embedded into complete  $SU(5)$  fundamental or anti-fundamental representation as  $\mathbf{5}_i = (\bar{D}'_i, L'_i)$  and  $\mathbf{\bar{5}}_i = (D'_i, \bar{L}'_i)$ ; we introduce  $N_5$  copies of  $\mathbf{5}$  and  $\mathbf{\bar{5}}$  with  $i = 1 \dots N_5$ .<sup>#2</sup> Then, the superpotential is given by

$$W = W_{\text{MSSM}} + M_V(\bar{D}'_i D'_i + \bar{L}'_i L'_i), \quad (1)$$

where  $W_{\text{MSSM}}$  is a superpotential of the MSSM and  $M_V$  is the common masses for vector-like matter fields.<sup>#3</sup> Hereafter,  $M_V (=M_{D'}=M_{L'})$  is taken to be  $\sim 1$  TeV, while  $N_5 = 3$  and 4, which are suggested by, for example, the diphoton excess observed by the LHC [13–30].<sup>#4</sup>

In order to see how the upper bound on the gluino mass is derived, it is instructive to see the leading one-loop correction to the Higgs mass. Assuming that the left- and right-handed stop masses are almost degenerate, the Higgs boson mass with the leading one-loop corrections in the decoupling limit is estimated as [2–6]

$$m_h^2 \simeq m_Z^2 \cos^2 2\beta + \frac{3}{4\pi^2} \frac{m_t^4}{v^2} \left[ \ln \frac{M_{\tilde{t}}^2}{m_t^2} + \frac{|X_t|^2}{M_{\tilde{t}}^2} \left( 1 - \frac{|X_t|^2}{12M_{\tilde{t}}^2} \right) \right], \quad (2)$$

where  $m_Z$  is the  $Z$ -boson mass,  $m_t$  is the top mass,  $M_{\tilde{t}}$  is the stop mass,  $v = 174.1$  GeV is the vacuum expectation value of the Higgs boson, and  $X_t = A_t - \mu/\tan\beta$  (with  $A_t$  being the trilinear coupling of stops normalized by the top Yukawa coupling constant  $y_t$ , and  $\mu$  being the Higgsino mass). Notice that the first term in the square bracket of Eq. (2) is the effect of the RG running of the quartic SM Higgs coupling constant from the mass scale of the SUSY particles to the electroweak scale, while the second one is the threshold correction at the mass scale of the SUSY particles. The Higgs mass becomes larger as  $M_{\tilde{t}}$  or  $X_t$  increases (as far as  $X_t \lesssim \sqrt{6}$ ). Thus, in order to realize the observed value of the Higgs mass,  $m_h \simeq 125$  GeV, there is an upper-bound on  $M_{\tilde{t}}$  and  $X_t$ . Importantly, the stop masses and the  $A_t$  parameter are enhanced with larger value of the gluino mass because of the RG runnings from a high

<sup>#2</sup>Our results are qualitatively unchanged even if the vector-like matters are embedded into other representations of  $SU(5)$ , as far as the parameter  $N_5$  is properly interpreted. For the case with  $N_{10}$  copies of  $\mathbf{\bar{10}}$  and  $\mathbf{10}$  representations, for example,  $N_5$  should be replaced by  $3N_{10}$ .

<sup>#3</sup>Due to the RG runnings, the SUSY invariant masses for  $D'$  and  $L'$  should differ even if they are unified at the GUT scale. Such an effect is, however, unimportant for our following discussion, and we neglect the mass difference among the extra matters.

<sup>#4</sup>The case of  $N_5 = 3$  is particularly interesting, since it may be embedded into an  $E_6$  GUT [35].

scale to the mass scale of SUSY particles. Consequently, with boundary conditions on the MSSM parameters given at a high scale, we obtain the upper bound on the gluino mass to have  $m_h \simeq 125$  GeV. Hereafter, we assume that the MSSM is valid up to the GUT scale  $M_{\text{GUT}} \sim 10^{16}$  GeV and derive such an upper bound.

Now we consider how the existence of the extra matters affects the upper bound on the gluino mass by using one-loop RG equations (RGEs), although two-loop RGEs are used for our numerical calculation in the next section. With  $N_5$  pairs of the vector-like multiplets, RGEs of gauge coupling constants at the one-loop level are

$$\frac{dg_i}{d \ln \mu_R} = \frac{b_i}{16\pi^2} g_i^3, \quad (3)$$

where  $\mu_R$  is a renormalization scale;  $g_1$ ,  $g_2$  and  $g_3$  are gauge coupling constants of  $U(1)_Y$  (in  $SU(5)$  GUT normalization),  $SU(2)_L$  and  $SU(3)_C$ , respectively. In addition,  $(b_1, b_2, b_3) = (33/5 + N_5, 1 + N_5, -3 + N_5)$ . For  $M_V \lesssim 1$  TeV,  $N_5 \lesssim 4$  needs to be satisfied under the condition that the coupling constants remain perturbative up to the GUT scale ( $\sim 10^{16}$  GeV). As one can see, for  $N_5 \gtrsim 3$ ,  $g_3$  is not asymptotically free. One-loop RGEs of gaugino masses are

$$\frac{dM_i}{d \ln \mu_R} = \frac{b_i}{8\pi^2} g_i^2 M_i, \quad (4)$$

where  $M_1$ ,  $M_2$  and  $M_3$  are the Bino, Wino and gluino mass, respectively. (Hereafter, we use the convention in which  $M_3$  is real and positive.) The ratio  $M_i/g_i^2$  is constant at the one-loop level and, with  $M_i$  at the mass scale of the SUSY particles being fixed, the gaugino masses at higher scale are more enhanced with larger value of  $N_5$ . In particular, for  $N_5 \gtrsim 3$ , the gluino mass, whose RG effects on  $A_t$  and the stop masses are important, become larger as the RG scale increases.<sup>#5</sup> In other words, even if  $|M_3|$  is large at the GUT scale, the low-energy value of  $|M_3|$  is small especially for  $N_5 = 4$  [26].

With the enhancement of the gluino mass, the RG effect on the  $A_t$  parameter becomes larger. This can be easily understood from the RGE of the  $A_t$  parameter; at the one-loop level,

$$\frac{dA_t}{d \ln \mu_R} = \frac{1}{16\pi} \left[ \frac{32}{3} g_3^2 M_3 + 6y_t^2 A_t + \cdots \right], \quad (5)$$

where we show only the terms depending on  $SU(3)_C$  gauge coupling constant or the top Yukawa coupling constant. One can see that the  $A_t$  parameter is generated by the RG effect using the gluino mass as a source, and the low-energy value of  $|A_t|$  is likely to become larger as  $|M_3|$  increases.

More quantitative discussion about the enhancement of the  $A_t$  parameter is also possible. Solving RGEs, the  $A_t$  parameter at the mass scale of the SUSY particles, denoted as  $m_S$ ,

---

<sup>#5</sup>For  $N_5 = 3$ , the one-loop beta-function vanishes and  $M_3$  is constant, but  $|M_3|$  becomes larger at the high energy scale due to two-loop effects.

can be parametrized as

$$A_t(m_{\mathcal{S}}) \simeq \begin{Bmatrix} -0.77 \\ -1.84 \\ -5.18 \end{Bmatrix} M_3(m_{\mathcal{S}}) + \begin{Bmatrix} 0.39 \\ 0.47 \\ 0.36 \end{Bmatrix} A_0, \quad (6)$$

where the numbers in the curly brackets are the coefficients for the cases of the MSSM (i.e.,  $N_5 = 0$ ),  $N_5 = 3$ , and  $N_5 = 4$ , from the top to the bottom, which are evaluated by using two-loop RGEs with  $m_{\mathcal{S}} = 3.5$  TeV, and  $A_0 \equiv A_t(M_{\text{inp}})$  with  $M_{\text{inp}}$  being the scale where the boundary conditions for the SUSY breaking parameters are set. (In our numerical calculations, we take  $M_{\text{inp}} = 10^{16}$  GeV.) In deriving Eq. (6) (as well as Eqs. (7) and (8)), we have taken  $\tan \beta = 10$ , and, for simplicity, we have assumed that (i) the gaugino masses obey the GUT relation, (ii) the SUSY breaking scalar masses are universal at  $M_{\text{inp}}$ , and (iii) all the trilinear scalar coupling constants are proportional to corresponding Yukawa coupling constant (with the proportionality factor  $A_0$ ) at  $M_{\text{inp}}$ . We can see that the coefficient of the  $M_3$  term becomes larger as  $N_5$  increases.

Similarly, we can discuss how the SUSY breaking stop mass parameters behave. Assuming the universality of the scalar masses at  $\mu_R = M_{\text{inp}}$ ,

$$m_{Q_3}^2(m_{\mathcal{S}}) \simeq \begin{Bmatrix} 0.68 \\ 2.36 \\ 9.89 \end{Bmatrix} M_3^2(m_{\mathcal{S}}) + \begin{Bmatrix} 0.05 \\ 0.21 \\ 1.06 \end{Bmatrix} M_3(m_{\mathcal{S}})A_0 + \begin{Bmatrix} -0.04 \\ -0.04 \\ -0.06 \end{Bmatrix} A_0^2 + \begin{Bmatrix} 0.66 \\ 0.61 \\ 0.53 \end{Bmatrix} \tilde{m}^2, \quad (7)$$

$$m_{\bar{U}_3}^2(m_{\mathcal{S}}) \simeq \begin{Bmatrix} 0.50 \\ 1.39 \\ 3.25 \end{Bmatrix} M_3^2(m_{\mathcal{S}}) + \begin{Bmatrix} 0.10 \\ 0.37 \\ 1.30 \end{Bmatrix} M_3(m_{\mathcal{S}})A_0 + \begin{Bmatrix} -0.08 \\ -0.08 \\ -0.05 \end{Bmatrix} A_0^2 + \begin{Bmatrix} 0.37 \\ 0.40 \\ 0.32 \end{Bmatrix} \tilde{m}^2, \quad (8)$$

where  $m_{Q_3}$  and  $m_{\bar{U}_3}$  are soft masses of the left-handed stop and right-handed stop, respectively, and  $\tilde{m}$  is the universal scalar mass. The coefficients of the  $M_3^2$  terms become significantly enhanced with larger value of  $N_5$ . Thus, with the increase of  $N_5$ , the stop masses becomes larger with fixed value of  $M_3^2(m_{\mathcal{S}})$ , as far as there is no accidental cancellation. In addition, Eqs. (6), (7) and (8) suggest that, for  $N_5 = 3$  and 4,  $A_t(m_{\mathcal{S}})$ ,  $m_{Q_3}^2(m_{\mathcal{S}})$ , and  $m_{\bar{U}_3}^2(m_{\mathcal{S}})$  are primarily determined by the gaugino mass if  $M_3$ ,  $A_0$ , and  $\tilde{m}$  are of the same size. We can see that, in such a case, the trilinear coupling constant  $A_t$  is more enhanced than the stop masses,  $m_{Q_3}$  and  $m_{\bar{U}_3}$ , which makes the threshold correction to the Higgs mass larger (see Eq. (2)).

Based on the above discussion, we have seen that the inclusion of the extra vector-like matter pushes up the Higgs mass for a fixed value of the gluino mass. Thus, the relevant value of the gluino mass realizing the observed Higgs mass becomes lower as the number of extra matter increases. In the next section, we will see that this is really the case, and derived the upper bound on the gluino mass with more detailed analysis of the RG effects.

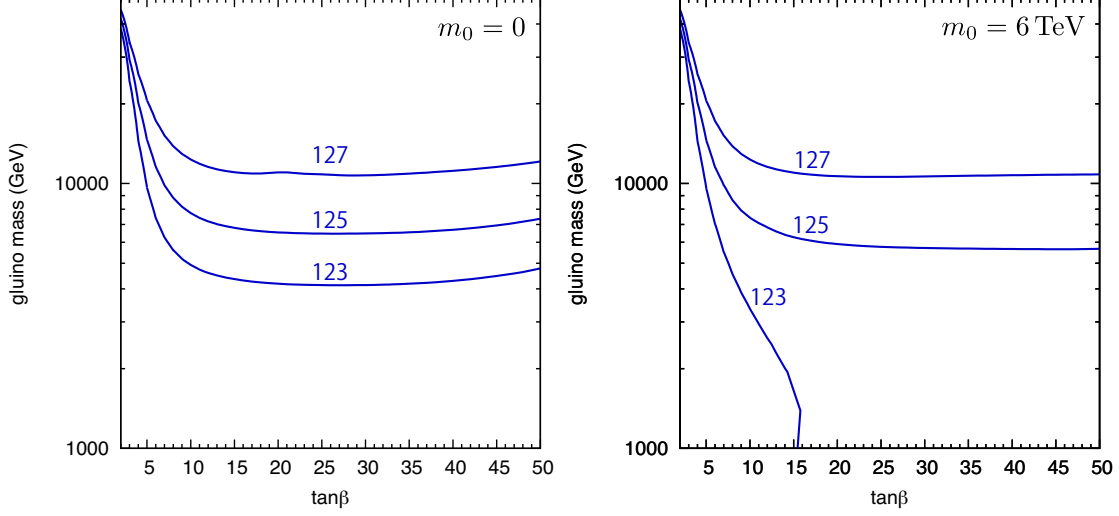


Figure 1: Contours of the Higgs boson mass in the MSSM. We take  $m_{H_u} = m_{H_d} = 0$ ,  $\mu > 0$  and  $A_0 = 0$ . The other scalar masses are set to be universal value  $m_0$ . The solid lines correspond to  $m_h = (127, 125, 123)$  GeV from top to bottom. Here,  $m_t(\text{pole}) = 173.34$  GeV and  $\alpha_s(m_Z) = 0.1185$ .

### 3 Numerical results

Now, we evaluate the upper-bound on the gluino mass by numerically solving two-loop RGEs. For our numerical calculation, we take  $m_t(\text{pole}) = 173.34$  GeV [37] and  $\alpha_s(m_Z) = 0.1185$  [38].

#### 3.1 The MSSM results

For the sake of comparison, we first show the upper-bound without vector-like multiplets, i.e. in the case of the MSSM. In the calculation, we take

$$M_1 = M_2 = M_3 = M_{1/2}, \quad m_{H_u} = m_{H_d} = 0, \quad A_0 = 0, \quad \mu > 0 \quad \text{at } M_{\text{inp}}, \quad (9)$$

with  $M_{\text{inp}} = 10^{16}$  GeV;  $A_0$  is the universal scalar trilinear coupling and  $m_{H_u}$  and  $m_{H_d}$  are the soft masses for the up-type and down-type Higgs, respectively. Scalar masses of MSSM matter multiplets ( $Q, \bar{U}, \bar{D}, L, \bar{E}$ ) are taken to be universal:

$$m_Q = m_{\bar{U}} = m_{\bar{D}} = m_L = m_{\bar{E}} = m_0 \quad \text{at } M_{\text{inp}}, \quad (10)$$

where we have omitted flavor indices. We choose  $m_{H_u} = m_{H_d} = 0$  rather than  $m_{H_u} = m_{H_d} = m_0$  in order to avoid the region with unsuccessful electroweak symmetry breaking (EWSB); if  $m_{H_u} = m_{H_d} = m_0$  and  $m_0 \gg M_{1/2}$ , the EWSB does not occurs [36]. Even if  $m_{H_u}$  and

$m_{H_d}$  are non-vanishing, the bound on the gluino mass is almost unchanged in most of the parameter space.<sup>#6</sup>

In Fig. 1, we show contours of the lightest Higgs boson mass,  $m_h$ , on  $\tan\beta$ - $m_{\tilde{g}}$  plane, where  $m_{\tilde{g}}$  is a physical gluino mass. The Higgs boson mass is computed by using **FeynHiggs** 2.11.3 [39–43]. The mass spectrum of the SUSY particles is calculated by using **SuSpect** 2.4.3 [44]. The blue solid lines show  $m_h$  of (127, 125, 123) GeV, from top to bottom. Although we expect the uncertainty in our calculation of the Higgs mass of a few GeV, we use the contour of  $m_h = 125$  GeV to discuss how the existence of the extra matter fields affects the upper bound on the gluino mass. Then, in the MSSM, the upper-bound on the gluino mass is as large as 7 TeV for  $\tan\beta > 10$  and  $A_0 = 0$ .<sup>#7</sup> Such a heavy gluino is hardly observed by the LHC experiment. In addition, if the gluino is so heavy, the squark masses are also expected to be so large via the RG effects unless there is an accidental cancellation. Thus, in the MSSM, the discovery of the colored SUSY particles is challenging unless the trilinear coupling of the stop is sizable at the boundary  $M_{\text{inp}}$ .

### 3.2 Gluino mass bound for $N_5 = 3$ and 4

Next we show the results for  $N_5 = 3$  and 4, for which the upper bound on the gluino mass is expected to be lower than the MSSM case. In addition to the boundary conditions Eqs. (9) and (10), we take the scalar masses for the vector-like multiplets to be universal:

$$m_{\bar{D}'_i} = m_{D'_i} = m_{\bar{L}'_i} = m_{L'_i} = m_0 \text{ at } M_{\text{inp}}. \quad (11)$$

The SUSY mass for the vector-like multiplets is taken to be  $M_V = 1$  TeV. SUSY mass spectra are calculated by solving two-loop RGEs with contributions from the vector-like multiplets and by including one-loop threshold corrections to the gauge coupling constants. These effects are included by modifying the **SuSpect** code. The one-loop threshold corrections from the vector-like multiplets are included by shifting the gauge couplings constants at the SUSY mass scale  $m_S$ :

$$\begin{aligned} g_1^{-2}(m_S) &\rightarrow g_1^{-2}(m_S) - \frac{N_5}{8\pi^2} \left[ \frac{2}{3} \ln \frac{m_S}{M_V} + \frac{1}{10} \ln \frac{m_S^2}{m_{L'_-} m_{L'_+}} + \frac{1}{15} \ln \frac{m_S^2}{m_{D'_-} m_{D'_+}} \right], \\ g_2^{-2}(m_S) &\rightarrow g_2^{-2}(m_S) - \frac{N_5}{8\pi^2} \left[ \frac{2}{3} \ln \frac{m_S}{M_V} + \frac{1}{6} \ln \frac{m_S^2}{m_{L'_-} m_{L'_+}} \right], \\ g_3^{-2}(m_S) &\rightarrow g_3^{-2}(m_S) - \frac{N_5}{8\pi^2} \left[ \frac{2}{3} \ln \frac{m_S}{M_V} + \frac{1}{6} \ln \frac{m_S^2}{m_{D'_-} m_{D'_+}} \right], \end{aligned} \quad (12)$$

where  $m_{L'_\pm}$  ( $m_{D'_\pm}$ ) are mass eigenvalues of the scalar components of  $L'_i$  and  $\bar{L}'_i$  ( $\bar{D}'_i$  and  $D'_i$ ).

<sup>#6</sup> The exception is the case where the EWSB occurs with the small  $\mu$ -parameter of a few hundred GeV, which will be discussed in Sec. 3.4.

<sup>#7</sup> In the region  $\tan\beta \simeq 1$ , the gluino mass bound is as high as  $10^{10}$  GeV [7].

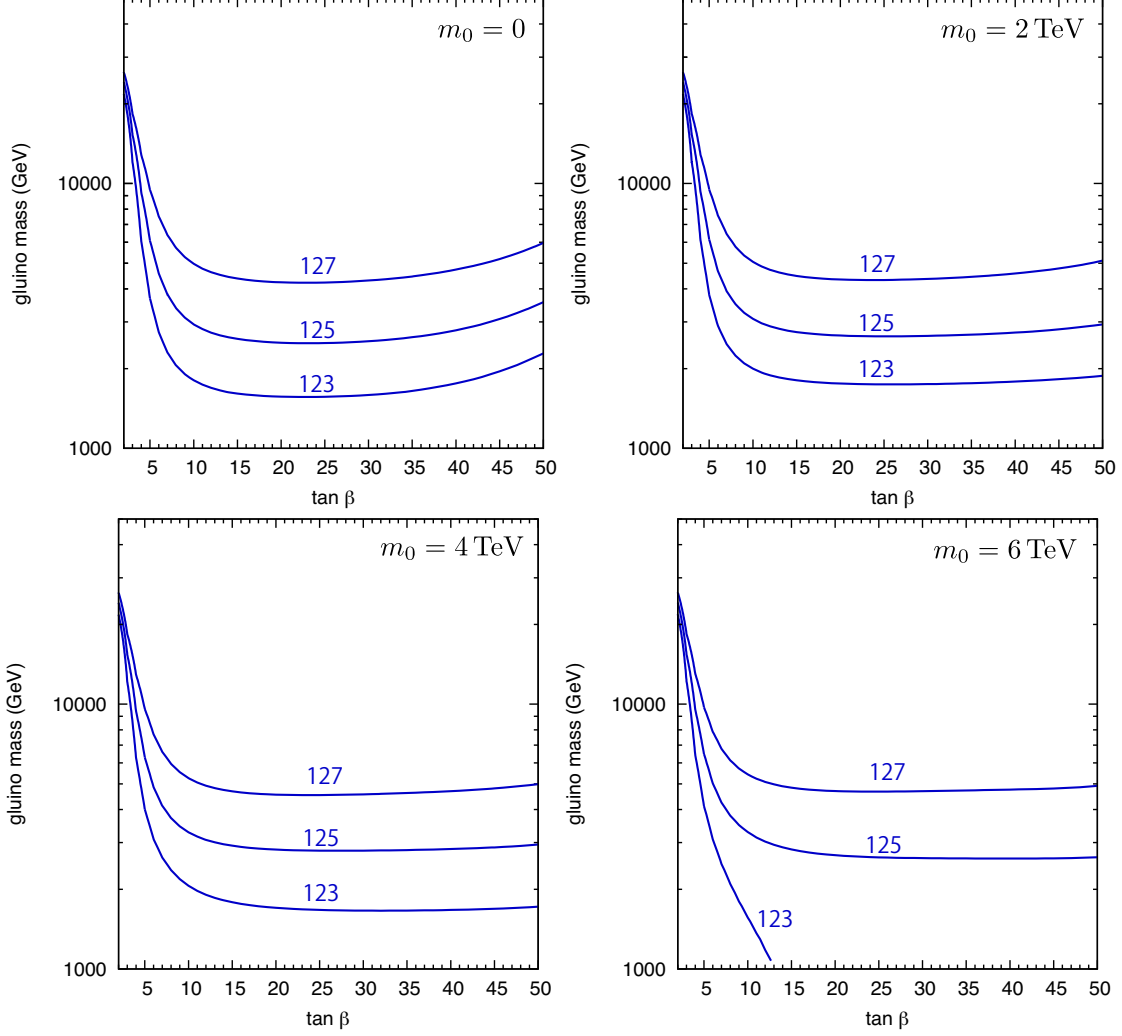


Figure 2: Contours of  $m_h$  on  $\tan\beta$ - $m_{\tilde{g}}$  plane for different  $m_0$ . We take  $N_5 = 3$  and  $M_V = 1$  TeV. The solid lines correspond to  $m_h = (127, 125, 123)$  GeV from top to bottom.

In Figs. 2 and 3, we show contours of  $m_h$  on  $\tan\beta$ - $m_{\tilde{g}}$  plane in the presence of vector-like multiplets for  $N_5 = 3$  and 4, respectively. The scalar mass  $m_0$  is taken to be  $m_0 = 0, 2, 4$ , and 6 TeV. In the case  $N_5 = 3$ ,  $m_h = 125$  GeV is realized with the gluino mass less than 3 TeV for large enough  $\tan\beta$  (i.e.,  $\tan\beta \gtrsim 10$ ). The gluino with such a mass is expected to be detectable with the high luminosity LHC [45]. In the case  $N_5 = 4$ , the upper-bound on the gluino mass is even smaller; the gluino should be lighter than  $\sim 2.5$  TeV for  $\tan\beta > 8$ .

In Figs. 4 and 5, we show contours of  $m_h$  (left),  $m_{\tilde{t}}$  and  $|A_t/m_{\tilde{t}}|^2$  (right) on  $m_0$ - $m_{\tilde{g}}$  plane, where  $m_{\tilde{t}} \equiv \sqrt{m_{Q_3} m_{\tilde{U}_3}}$ . In Fig. 4, we take  $N_5 = 3$  and  $\tan\beta = 25$ . The gluino mass is smaller than 2.8 TeV for  $m_h = 125$  GeV, because of sizable  $|A_t/m_{\tilde{t}}|^2$ , which enhance the threshold correction to the Higgs mass, or relatively large  $m_{\tilde{t}} > 5$  TeV. In Fig. 5, we take  $N_5 = 4$  and



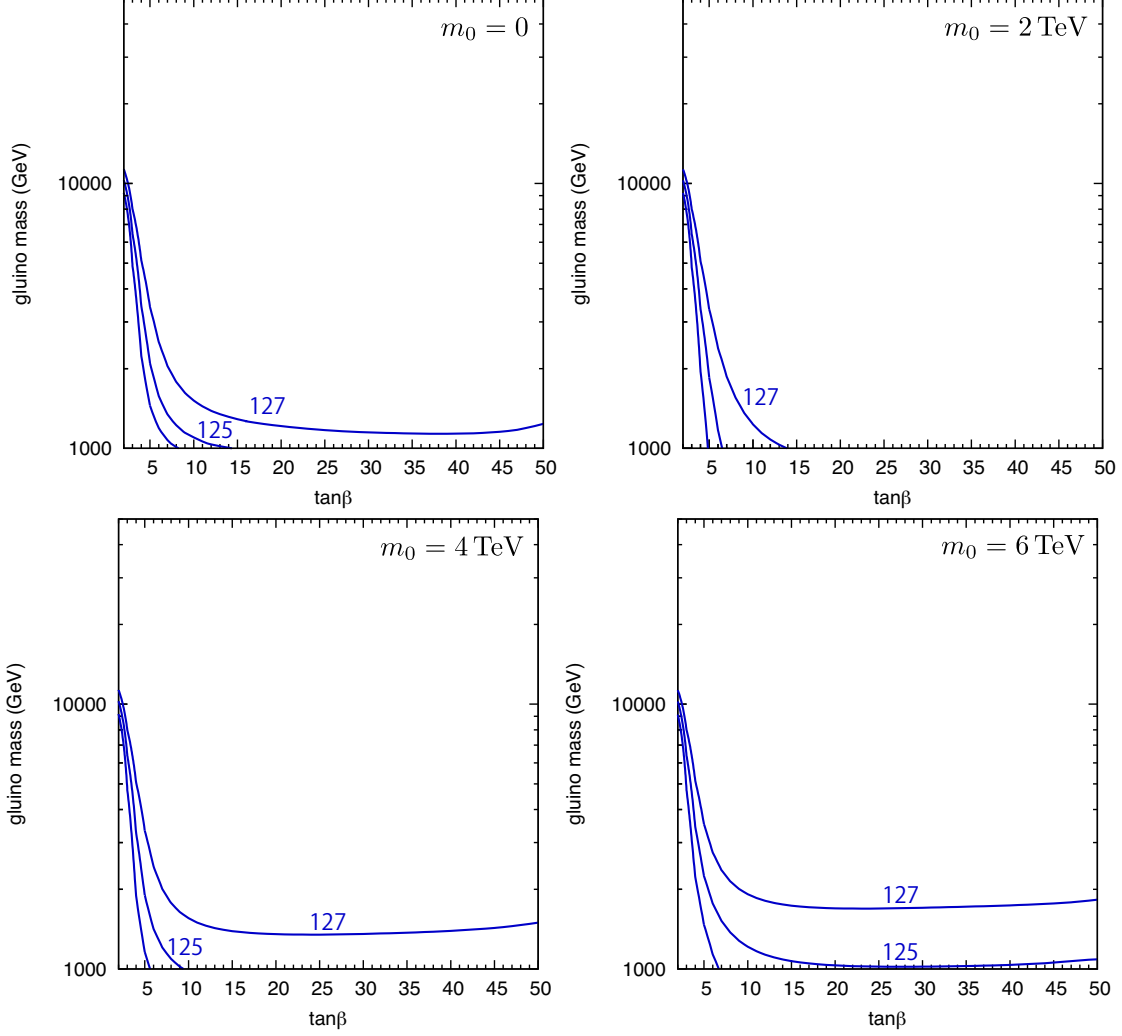


Figure 3: Contours of  $m_h$  on  $\tan\beta$ - $m_{\tilde{g}}$  plane for different  $m_0$ . The solid lines correspond to  $m_h = (127, 125, 123)$  GeV from top to bottom. Here,  $N_5 = 4$ .

$\tan\beta = 5$ . The gluino mass is smaller than 2.5 TeV due to the large  $|A_t/m_t|^2$ , which is as large as 4–5 for  $m_0 < 2$  TeV. We also comment on the  $m_0$ -dependence of the bound on the gluino mass. When  $m_0$  is relatively small, the stop masses are determined mostly by the gluino mass via the RG effects. In such a case, the upper bound on the gluino mass is insensitive to  $m_0$ . On the contrary, when  $m_0$  is large, the stop masses becomes sensitive to  $m_0$ ; in such a case, with the increase of  $m_0$ , the upper bound on the gluino mass becomes lower.

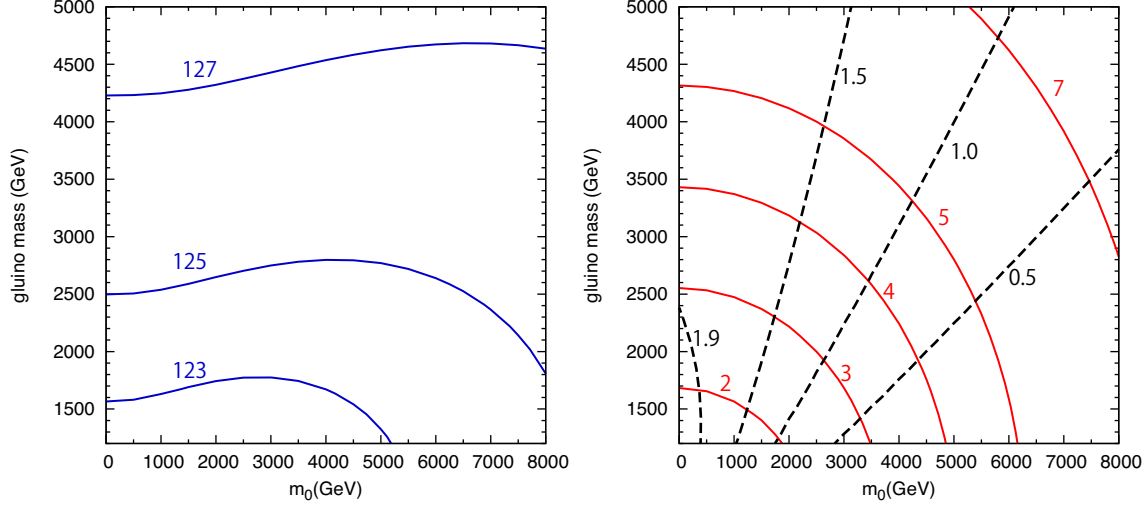


Figure 4: Contours of  $m_h$  in unit of GeV (left),  $m_t$  and  $|A_t^2|/m_t^2$  (right) for  $N_5 = 3$ . The stop mass in unit of TeV and  $|A_t^2|/m_t^2$  are shown in the red solid line and black dashed line, respectively. Here,  $\tan \beta = 25$ .

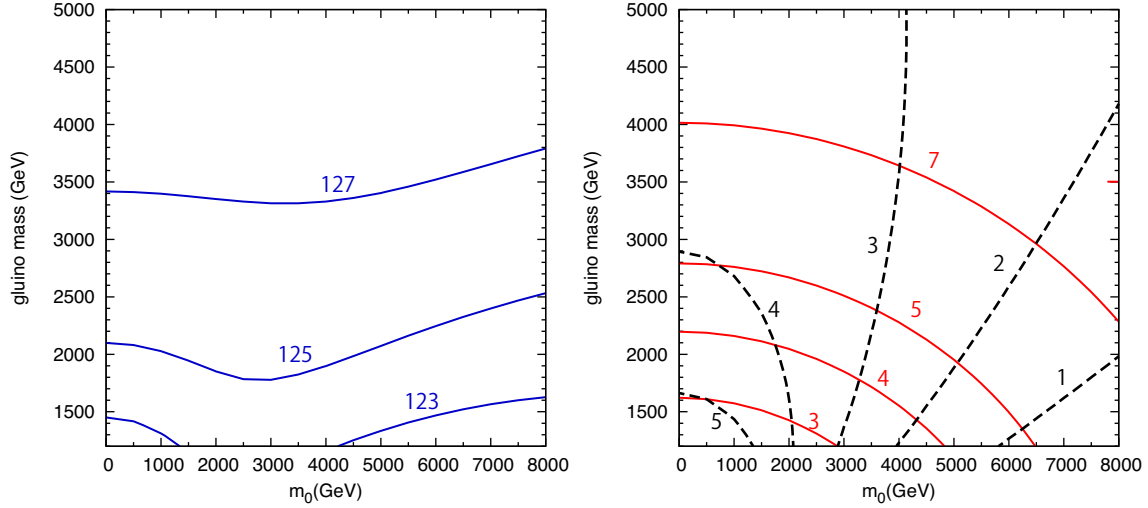


Figure 5: Contours of the  $m_h$  (left),  $m_t$  and  $|A_t^2|/m_t^2$  (right) for  $N_5 = 4$ . Here,  $\tan \beta = 5$ .

### 3.3 The effect of the bare $A$ -term

Let us discuss the effects of the non-zero  $A_0$ . The bare  $A$ -parameter,  $A_0$ , contributes to  $A_t$  destructively (constructively) if  $M_{1/2}$  and  $A_0$  have same (opposite) signs (see Eq. (6)). Thus, with taking negative  $A_0$ ,  $A_t(m_S)$  increase, and the gluino mass which realizes  $m_h \simeq 125$  GeV becomes smaller. On the contrary, when  $A_0$  is positive and large,  $A_t(m_S)$  becomes suppressed

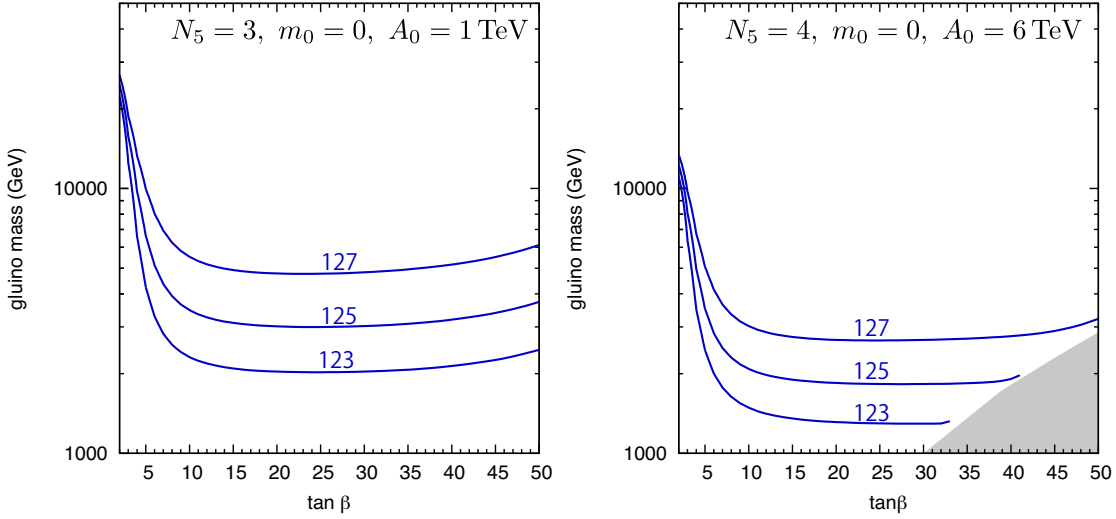


Figure 6: The Higgs boson mass with non-zero  $A_0$ . In the left (right) panel, we take  $N_5 = 3(4)$  and  $A_0 = 1(6)$  TeV. The gray region in the right panel is excluded due to the tachyonic stau.

so that the upper bound on the gluino mass can become higher. Notice that the effects of non-vanishing  $A_0$  on the stop masses are not so significant unless  $|A_0|$  is very large (see Eqs. (7) and (8)).

In Fig. 6, we show the upper bound on the gluino mass, taking non-vanishing  $A_0 > 0$ . In the case  $N_5 = 3$  and  $A_0 = 1$  TeV, the gluino mass bound slightly increases compared to the case of  $A_0 = 0$ . The gluino mass bound remains  $\sim 3$  TeV for  $15 < \tan \beta < 35$ . In the case  $N_5 = 4$ , even if  $A_0 = 6$  TeV, the gluino mass bound is as small as 2 TeV for  $\tan \beta > 10$ .

### 3.4 Implications of dark matter

In the simple set up discussed above, the Bino-like neutralino is the lightest SUSY particle (LSP) and its relic density is much larger than the observed dark matter density,  $\Omega_{\text{CDM}} h^2 \simeq 0.12$  [55]; in such a case, the thermal relic lightest neutralino is not a viable dark matter candidate. However, with a slight modification, the lightest neutralino can be dark matter without much affecting the gluino mass bound. Below, we discuss several possibilities.

- Higgsino dark matter

If  $m_{H_u}^2(M_{\text{inp}})$  is tuned, the EWSB can occur with  $\mu \sim m_Z$ . In such a case, the Higgsino of a few hundred GeV can be the LSP. The thermal relic abundance of such Higgsino LSP is smaller by  $\sim 1/10$  compared to the observed dark matter abundance. However, the observed dark matter abundance can be realized with non-thermal productions [46–48]. With smaller value of the Higgsino mass, the Higgs boson mass is raised about

Table 1: Mass spectra in sample points. We take  $A_0 = 0$  and  $M_{\text{inp}} = 10^{16} \text{ GeV}$ . Here,  $A_t$  shown in the table is the generated  $A$ -term at  $m_S$ .

Parameters	Point I	Point II	Point III	Point IV	Point V
$N_5$	3	3	4	4	3
$M_3 \text{ (GeV)}$	3000	3540	6900	6300	3400
$M_1/M_3$	1	1.0	0.83	1	0.7
$M_2/M_3$	1	0.61	0.62	1	1
$m_0 \text{ (GeV)}$	0	0	0	4000	0
$m_{H_{u,d}}/1 \text{ TeV}$	3.441	0	0	6.392	0
$\tan \beta$	10	25	6	5	32.9
$\mu \text{ (GeV)}$	229	3410	5270	194	3210
$A_t \text{ (GeV)}$	-4030	-4450	-7050	-6720	-4480
Particles	Mass (GeV)	Mass (GeV)	Mass (GeV)	Mass (GeV)	Mass (GeV)
$\tilde{g}$	2470	2970	1890	1760	2840
$\tilde{q}$	3670–3890	4340–4400	5360–5370	5900–6070	4150–4410
$\tilde{t}_{2,1}$	3220, 2130	3780, 3250	4390, 3130	4760, 2560	3720, 2940
$\tilde{\chi}_{2,1}^\pm$	942, 232	3410, 669	5240, 537	884, 196	3210, 1110
$\tilde{\chi}_4^0$	942	3410	5240	884	3210
$\tilde{\chi}_3^0$	537	3410	5240	590	3210
$\tilde{\chi}_2^0$	238	669	537	203	1110
$\tilde{\chi}_1^0$	227	645	510	191	420
$\tilde{e}_{L,R}(\tilde{\mu}_{L,R})$	1510, 911	1140, 1080	1630, 1420	4580, 4260	1700, 715
$\tilde{\tau}_{2,1}$	1500, 860	1150, 983	1630, 1420	4580, 4250	1650, 423
$H^\pm$	3730	3290	5590	6910	3040
$h_{\text{SM-like}}$	125.2	125.2	126.3	125.6	125.2

1 GeV for the fixed gluino mass, and hence the bound on the gluino mass becomes stronger.

- Bino-Wino coannihilation

If we relax the GUT relation among the gaugino masses, the mass difference between the Bino and Wino can be small and the thermal relic abundance of the lightest neutralino can be consistent with the observed dark matter density, because of the Bino-Wino coannihilation with  $M_2/M_1(M_{\text{inp}}) \sim 0.5-0.7$ . This reduces the Higgs boson mass only slightly. Thus, the upper-bound on the gluino mass is almost unchanged.

- Wino dark matter

If the mass ratio of  $M_2/M_1$  is even smaller than the previous case, the lightest neutralino can be (almost) Wino-like. Although, the thermal relic abundance of the Wino-like

neutralino is too small as in the case of the Higgsino dark matter, with non-thermal production, the relic abundance can be consistent with the observed dark matter relic. The constraint from gamma rays from dwarf spheroidal galaxies gives a lower-bound on the Wino mass,  $M_2(m_{\mathcal{S}}) \gtrsim 320 \text{ GeV}$  [49]. The LHC may discover/exclude the Wino LSP with the mass up to  $\sim 500 \text{ GeV}$  through electroweak productions [50].

- Bino-stau coannihilation

For  $N_5 = 3$  and large  $\tan \beta$ , the mass difference between stau and neutralino becomes small and the relic abundance of the neutralino is reduced due to the Bino-stau coannihilation [51, 52]. Smaller  $M_1$  at  $M_{\text{inp}}$  helps to reduce the mass difference.

In the Table 1, we show sample points where the relic density of the lightest neutralino is consistent with the observed dark matter density. We calculate the thermal relic density using `MicrOMEGAs 4.1.7` package [53, 54]. In the points **I**, **II** and **V** (**III** and **IV**), we take  $N_5 = 3$  ( $N_5 = 4$ ). In the points **II**, **III** and **V**, the GUT relation among the gaugino masses is relaxed. In **II** and **III** (**V**), because of the Bino-Wino coannihilation (stau coannihilation), the thermal relic abundance of the lightest neutralino becomes constraint from the observed dark matter abundance. In the points **I** and **IV**, the Higgsino like neutralino is the LSP, and the thermal relic abundance smaller by about 1/10 compared to the observed dark matter abundance. With non-thermal productions, this Higgsino-like neutralino can be candidate for a dark matter.<sup>#8</sup>

## 4 Conclusion and discussion

We have investigated the upper-bound on the gluino mass for  $m_h \simeq 125 \text{ GeV}$ , with 3 or 4 copies of the vector-like multiples around TeV, transforming **5** and  **$\bar{5}$**  representations in  $SU(5)$  GUT gauge group. We have shown that with these vector-like multiplets, the upper-bound on the gluino mass is significantly reduced compared to that of the MSSM. The significant reduction originates from the fact that the radiatively generated trilinear coupling of stops as well as the stop masses is enhanced for the fixed gluino mass at the low-energy scale. In both cases of  $N_5 = 3$  and 4, the gluino mass is less than 3 TeV in a wide range of parameter space, and the gluino is likely to be discovered at the LHC Run-2 or the high luminosity LHC.

The presence of 3 or 4 copies of the vector-like multiplets at around TeV is suggested by the recently observed 750 GeV diphoton excess at the LHC. With the vector-like extra matters as well as a gauge singlet field  $\Phi$ , the diphoton excess can be explained by the production and the diphoton decay of  $\Phi$ . In such a scenario, the singlet field  $\Phi$  should couple to the vector-like matters **5** and  **$\bar{5}$**  with Yukawa couplings of  $\sim 1$  in the superpotential. In order to make the mass of  $\Phi$  as well as those of the vector-like matters being close to the TeV

---

<sup>#8</sup>One might worry about the constraint from the direct detection experiments of dark matter. However, the neutralino-nucleon scattering cross section is suppressed when  $M_1$  is large.

scale, it is attractive if they have the same origin, i.e., the vacuum expectation value of  $\Phi$ ,  $\langle\Phi\rangle \simeq \text{TeV}$ . Even if the SUSY breaking soft mass squared of  $\Phi$ , denoted as  $m_\Phi^2$ , is positive at a high scale, it can be driven to negative via the RG effect because  $\Phi$  is expected to have a strong Yukawa interaction with the vector-like matters. With such a negative mass squared parameter,  $\langle\Phi\rangle$  can become non-vanishing. However, we found that  $|m_\Phi^2|$  is typically much larger than  $(1\text{ TeV})^2$  in the parameter region of our interest. Thus, we need a tuning to obtain the correct size of  $\langle\Phi\rangle$  by introducing supersymmetric mass parameters for  $\Phi$ .<sup>#9</sup> the required tuning is typically  $\mathcal{O}(1)$  percent level.

## Acknowledgments

This work is supported by JSPS KAKENHI Grant Numbers JP26400239 (T.M.), JP26104009 (T.T.Y.), JP26287039 (T.T.Y.), JP16H02176 (T.T.Y.), JP15H05889 (N.Y.) and JP15K21733 (N.Y.); and by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan (T.M and T.T.Y.).

## References

- [1] G. Aad *et al.* [ATLAS and CMS Collaborations], Phys. Rev. Lett. **114**, 191803 (2015) [arXiv:1503.07589 [hep-ex]].
- [2] Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. **85**, 1 (1991).
- [3] Y. Okada, M. Yamaguchi and T. Yanagida, Phys. Lett. B **262**, 54 (1991).
- [4] J. R. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B **257**, 83 (1991).
- [5] J. R. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B **262**, 477 (1991).
- [6] H. E. Haber and R. Hempfling, Phys. Rev. Lett. **66**, 1815 (1991).
- [7] E. Bagnaschi, G. F. Giudice, P. Slavich and A. Strumia, JHEP **1409**, 092 (2014) [arXiv:1407.4081 [hep-ph]].
- [8] The ATLAS collaboration, ATLAS-CONF-2015-081.
- [9] The ATLAS collaboration, ATLAS-CONF-2016-018.
- [10] CMS Collaboration, CMS-PAS-EXO-15-004.
- [11] CMS Collaboration [CMS Collaboration], CMS-PAS-EXO-16-018.

---

<sup>#9</sup> Even in a case that the masses of the vector-like matters do not originate from the non-zero  $\langle\Phi\rangle$ , in order to make a mass of a scalar component of  $\Phi$  to be 750 GeV, one generally requires a tuning due to the negative  $m_\Phi^2$ .

- [12] A. Angelescu, A. Djouadi and G. Moreau, Phys. Lett. B **756**, 126 (2016) [arXiv:1512.04921 [hep-ph]].
- [13] D. Buttazzo, A. Greljo and D. Marzocca, Eur. Phys. J. C **76**, no. 3, 116 (2016) [arXiv:1512.04929 [hep-ph]].
- [14] R. Franceschini *et al.*, JHEP **1603**, 144 (2016) [arXiv:1512.04933 [hep-ph]].
- [15] J. Ellis, S. A. R. Ellis, J. Quevillon, V. Sanz and T. You, JHEP **1603**, 176 (2016) [arXiv:1512.05327 [hep-ph]].
- [16] K. M. Patel and P. Sharma, Phys. Lett. B **757**, 282 (2016) [arXiv:1512.07468 [hep-ph]].
- [17] N. Craig, P. Draper, C. Kilic and S. Thomas, arXiv:1512.07733 [hep-ph].
- [18] L. J. Hall, K. Harigaya and Y. Nomura, JHEP **1603**, 017 (2016) [arXiv:1512.07904 [hep-ph]].
- [19] Y. L. Tang and S. h. Zhu, arXiv:1512.08323 [hep-ph].
- [20] F. Wang, W. Wang, L. Wu, J. M. Yang and M. Zhang, arXiv:1512.08434 [hep-ph].
- [21] E. Palti, Nucl. Phys. B **907**, 597 (2016) [arXiv:1601.00285 [hep-ph]].
- [22] W. Chao, arXiv:1601.00633 [hep-ph].
- [23] B. Dutta, Y. Gao, T. Ghosh, I. Gogoladze, T. Li, Q. Shafi and J. W. Walker, arXiv:1601.00866 [hep-ph].
- [24] A. Djouadi, J. Ellis, R. Godbole and J. Quevillon, JHEP **1603**, 205 (2016) [arXiv:1601.03696 [hep-ph]].
- [25] S. F. King and R. Nevzorov, JHEP **1603**, 139 (2016) [arXiv:1601.07242 [hep-ph]].
- [26] C. Han, T. T. Yanagida and N. Yokozaki, Phys. Rev. D **93**, no. 5, 055025 (2016) [arXiv:1602.04204 [hep-ph]].
- [27] R. Barbieri, D. Buttazzo, L. J. Hall and D. Marzocca, arXiv:1603.00718 [hep-ph].
- [28] H. P. Nilles and M. W. Winkler, JHEP **1605**, 182 (2016) [arXiv:1604.03598 [hep-ph]].
- [29] L. J. Hall, K. Harigaya and Y. Nomura, arXiv:1605.03585 [hep-ph].
- [30] T. Cohen, G. D. Kribs, A. E. Nelson and B. Ostdiek, arXiv:1605.04308 [hep-ph].
- [31] K. J. Bae, M. Endo, K. Hamaguchi and T. Moroi, Phys. Lett. B **757** (2016) 493 [arXiv:1602.03653 [hep-ph]].

- [32] K. Kurosawa, N. Maru and T. Yanagida, Phys. Lett. B **512**, 203 (2001) [hep-ph/0105136].
- [33] M. Asano, T. Moroi, R. Sato and T. T. Yanagida, Phys. Lett. B **705**, 337 (2011) [arXiv:1108.2402 [hep-ph]].
- [34] T. Moroi, T. T. Yanagida and N. Yokozaki, Phys. Lett. B **719**, 148 (2013) [arXiv:1211.4676 [hep-ph]].
- [35] See, for a recent work, Y. Cai, J. D. Clarke, R. R. Volkas and T. T. Yanagida, arXiv:1605.02743 [hep-ph].
- [36] J. L. Feng, K. T. Matchev and T. Moroi, Phys. Rev. Lett. **84**, 2322 (2000) doi:10.1103/PhysRevLett.84.2322 [hep-ph/9908309].
- [37] [ATLAS and CDF and CMS and D0 Collaborations], arXiv:1403.4427 [hep-ex].
- [38] K.A. Olive et al. (Particle Data Group), Chin. Phys. C, **38**, 090001 (2014).
- [39] S. Heinemeyer, W. Hollik and G. Weiglein, Comput. Phys. Commun. **124**, 76 (2000) [hep-ph/9812320].
- [40] S. Heinemeyer, W. Hollik and G. Weiglein, Eur. Phys. J. C **9**, 343 (1999) [hep-ph/9812472].
- [41] G. Degrandi, S. Heinemeyer, W. Hollik, P. Slavich and G. Weiglein, Eur. Phys. J. C **28**, 133 (2003) [hep-ph/0212020];
- [42] M. Frank, T. Hahn, S. Heinemeyer, W. Hollik, H. Rzehak and G. Weiglein, JHEP **0702**, 047 (2007) [hep-ph/0611326].
- [43] T. Hahn, S. Heinemeyer, W. Hollik, H. Rzehak and G. Weiglein, Phys. Rev. Lett. **112**, no. 14, 141801 (2014) [arXiv:1312.4937 [hep-ph]].
- [44] A. Djouadi, J. L. Kneur and G. Moultaka, Comput. Phys. Commun. **176**, 426 (2007) [hep-ph/0211331].
- [45] The ATLAS collaboration, ATL-PHYS-PUB-2014-010.
- [46] G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998) [hep-ph/9810442].
- [47] T. Moroi and L. Randall, Nucl. Phys. B **570**, 455 (2000) [hep-ph/9906527].
- [48] G. B. Gelmini and P. Gondolo, Phys. Rev. D **74**, 023510 (2006) [hep-ph/0602230].
- [49] B. Bhattacharjee, M. Ibe, K. Ichikawa, S. Matsumoto and K. Nishiyama, JHEP **1407**, 080 (2014) [arXiv:1405.4914 [hep-ph]].



- [50] T.Yamanaka, Progress in Particle Physic (2013), <http://www2.yukawa.kyoto-u.ac.jp/~ppp/PPP2013/slides/YamanakaT.pdf> .
- [51] J. R. Ellis, T. Falk and K. A. Olive, Phys. Lett. B **444**, 367 (1998) [hep-ph/9810360].
- [52] J. R. Ellis, T. Falk, K. A. Olive and M. Srednicki, Astropart. Phys. **13**, 181 (2000) Erratum: [Astropart. Phys. **15**, 413 (2001)] [hep-ph/9905481].
- [53] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, Comput. Phys. Commun. **149**, 103 (2002) [hep-ph/0112278].
- [54] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, Comput. Phys. Commun. **174**, 577 (2006) [hep-ph/0405253].
- [55] R. Adam *et al.* [Planck Collaboration], arXiv:1502.01582 [astro-ph.CO].